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STUDY OF POTENTIALLY DANGEROUS EVENTS DURING HEAVY-ION COLLISIONS AT THE LHC: REPORT OF THE LHC SAFETY STUDY GROUP

J.-P. Blaizot
CEA/Saclay-Orme des Merisiers, Gif-sur-Yvette, France

J. Iliopoulos
École Normale Supérieure, Paris, France

J. Madsen,
University of Aarhus, Århus, Denmark

G.G. Ross,
University of Oxford, Oxford, UK

P. Sonderegger,
CERN, Geneva, Switzerland

H.-J. Specht,
University of Heidelberg, Heidelberg, Germany

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Abstract

We review the possibility of producing dangerous objects during heavy-ion collisions at the Large Hadron Collider. We consider all such objects that have been theoretically envisaged, such as negatively charged strangelets, gravitational black holes, and magnetic monopoles. We find no basis for any conceivable threat.

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Contents

1 INTRODUCTION

The desire to explore shorter and shorter distances has driven particle physicists to design and build accelerators with higher and higher energies. The Large Hadron Collider (LHC), at present under construction at CERN, will be the highest energy accelerator ever built. Many new and exciting phenomena are expected to occur as a result of these very high energy collisions. It is hoped that some of them will be unpredicted and will point to new directions in our understanding of the structure of matter. At the same time it is legitimate to wonder whether any of these new phenomena may be potentially dangerous. Such concerns were expressed for the first time before the commissioning and operation of the Brookhaven Relativistic Heavy Ion Collider (RHIC). A committee was set up by Prof. John Marburger, Director of Brookhaven National Laboratory, to review the case and a report was published in September 1999 [1]. Since the LHC will reach even higher energies in accelerating heavy ions, Prof. Luciano Maiani, CERN's Director-General, has charged this committee to conduct a new and comprehensive study in view of all new developments, both in theoretical studies and experimental results. This report summarizes our conclusions. Two general classes of phenomenon have been considered: negatively charged strangelets and exotic objects associated with gravitational interactions. In Sections 2 and 3 these two questions are addressed, with emphasis on new developments since the Brookhaven report.

On the first issue, that of strangelet production, we have considered new theoretical studies on the stability of negatively charged strangelets [2], as well as the lessons learnt from all data, including those from RHIC, on the production of complex nuclei.

No rigorous proofs of the existence, let alone the stability, of strange quark matter can be given, and all quantitative studies so far have used simple hadronic models like the MIT bag model. Lattice calculations, in principle using only fundamental QCD properties, are still far from describing complex systems with a large number of fermions. Nevertheless, a recent study shows that negatively charged strangelets cannot be stable for all values of baryon number A. Therefore, even if such an object is produced for sufficiently low A and starts growing, it will soon reach the instability region and stop growing. The detailed calculation again uses the MIT bag model, but the argument is essentially based on fundamental physical principles and we believe it to be robust. It is briefly presented in Section 2.1. Further details can be found in the original publication [2]. This result is sufficient to show that strangelets do not offer a realistic threat. However, in order to explore every aspect of the disaster scenario, however implausible, we also consider whether significant strangelet production is conceivable at the LHC, ignoring the likely impossibility of strangelet growth.

We remind the reader that for a potentially dangerous object to be produced at the LHC, three conditions are necessary. The production cross-section must be sharply peaked in the central rapidity region, otherwise the existing astrophysical data, such as cosmic ray collisions on the surface of the moon, set limits beyond the reach of any accelerator. Even then, it must have a threshold on the baryon number of the colliding heavy ions, since, as shown in Ref. [3], head-on collisions of high-energy heavy ions occur in outer space for baryon numbers such as iron. And finally, it must have an energy threshold to make it accessible only to the LHC and not to lower energy accelerators.

In Section 2.2 we review briefly the evidence and the limits from cosmic-ray data. We know of no consistent model of hadron production which satisfies all these conditions, so in the spirit of the devil's advocate we follow a purely phenomenological approach. Strangelets, if they exist at all, are hadronic systems made out of quarks and any model for their production should be first tested against the existing data on the production of nuclei. Indeed, in recent years a fair amount of such data has been accumulated from the Alternating Gradient Synchrotron (AGS) at BNL, the Super Proton Synchrotron (SPS) at CERN, and now RHIC, providing us with quite reliable phenomenological constraints on the underlying production mechanisms. In Section 2.3 we attempt to present a comprehensive review of this topic. The main lesson is that coalescence models describe these processes in the entire energy range and for all baryon numbers reached so far with sufficient accuracy to allow reliable extrapolations. Coalescence is based on a simple statistical argument and is summarized in Section 2.3.2. We go beyond

the estimations presented in the Brookhaven report by including the correct scaling with beam energy. It is remarkable that no additional mechanism is necessary to fit the data. This point is also developed in Section 2.3.2, where we consider in addition thermal and distillation models. Our conclusions are supported by the first results obtained at RHIC, which we have included in our estimations. We note that in the years to come before LHC operation this set of new data will be largely expanded; we intend to monitor and adjust our conclusions in the light of such data. In fact, in Section 2.3 we argue that the LHC will not be more efficient in producing strangelets than RHIC.

The second question we address is the possibility of creating dangerous objects associated with gravitational interactions. A similar question was also discussed in the RHIC report with the conclusion, expected by ordinary dimensional analysis, that such gravitational effects are suppressed by inverse powers of the Planck mass M_{P} and are, therefore, negligible. Recently, however, there have been suggestions that M_{P} is not the right parameter to use in the analysis because it does not determine the fundamental scale of the theory. These models contain extra compact space dimensions [4], whose size may be much larger than ${M_{\mathrm{P}}}^{-1}$, in fact as large as the inverse of a few TeV. This opens the exciting possibility of observing the effects of these extra dimensions at the LHC, but also requires a new examination of potential hazards. We present our estimations in Section 3 with the conclusion that dangerous objects, like growing black holes, are still far beyond the reach of the LHC, essentially because only extremely massive ones are stable.

Before closing this introduction we wish to make a general remark: All estimates concerning production probabilities and subsequent properties of various objects at the LHC necessarily involve certain theoretical assumptions. Some, for example the invariance of physical laws under space and time translations, are so general that they do not need to be explicitly stated. Others are based on extrapolations of known properties of hadronic systems. They will be explained in the following sections whenever they are used. Here we want only to emphasize that no estimates are absolutely assumption-free.

2 STRANGELETS

2.1 Strangelet properties

Quark matter containing roughly equal numbers of up, down, and strange quarks could be metastable or even absolutely stable (i.e. the ground state of hadronic matter), in spite of the fact that ordinary nuclei do not decay into such strange quark matter. Stability of nuclei demonstrates that two-flavour quark matter is not stable, but the addition of an extra Fermi sea for strange quarks lowers the energy relative to a two-flavour system for fixed baryon number, as long as the s-quark is not too heavy. Decay of ordinary nuclei into this state is prevented by the necessity to create s-quarks by weak interactions.

No first-principles theory of strange quark matter or strangelets exists at the moment. This would require a breakthrough in the application of QCD to finite-density systems. Instead, strange quark matter has been studied within phenomenological models, in particular the MIT bag model where confinement is described in terms of a bag pressure (the 'bag' constant B), gluon exchange in terms of a strong 'fine-structure' constant α_{S}, and quarks are treated as Fermi gases with current quark masses $m_{\mathrm{u}} \approx m_{\mathrm{d}} \approx 0$, $m_{\mathrm{s}} > 0$.

The physical properties of strangelets are described in detail in the Brookhaven report [1]. Here we summarize the main aspects of relevance for potential disaster scenarios [5], including some comments and new information that goes beyond the Brookhaven report.

(1) Bulk strange quark matter may be absolutely stable, but only for a somewhat narrow parameter window. 'Best guesses' for the relevant parameter values tend to disfavour strange quark matter stability. For example, a naive use of bag models to describe the quark–hadron phase transition at zero chemical potential would imply that the bag constant B and the phase transition temperature T_{C} are related as $T_{\mathrm{C}} \approx 0.65$–$0.7B^{1/4}$, leading to $B^{1/4} > 210$ MeV according to lattice

calculations, whereas strange quark matter stability requires $B^{1/4} < 165$ MeV. On the other hand, the MIT bag model fits ordinary hadron properties with $B^{1/4} = 145$ MeV (but these fits involve several other adjustable parameters). It is fair to say, though, that it is not obvious that phenomenological bag models should describe quark matter at high and low chemical potential using one value for the bag constant.

(2) Even if bulk strange quark matter is stable, finite size effects (surface tension and curvature) significantly destabilize strangelets with low baryon number, for typical parameters, adding 50 MeV per baryon for $A = 20$ and 85 MeV per baryon for $A = 10$. Since only very low-mass objects are potentially formed in heavy-ion collisions, this destabilization makes strangelet formation in colliders much less likely.

(3) Strangelets formed in heavy-ion collisions would initially have a significant entropy (temperature). Hot strangelets are much less stable than cold strangelets, and the time-scale for evaporation would presumably be very short (but difficult to calculate in detail). Very little is known about the decay channels and time-scales even at zero temperature. Calculations exist for the weak conversion time-scales for quarks inside strange quark matter, but coupling these to particle emission rates from a strangelet is essentially impossible at present.

(4) Potential disaster scenarios only occur if the total quark charge of strange quark matter is negative, so that nuclei are electrostatically attracted rather than repelled. Most calculations of strange quark matter properties find the charge to be positive. This is because the abundance of s-quarks is somewhat reduced relative to u- and d-quarks due to the higher mass of the s-quark. But one-gluon exchange interactions may change this for high values of the strong interaction fine-structure constant α_S because they are attractive for the massive s-quark and repulsive for massless u- and d-quarks. This increases the relative abundance of negatively charged s-quarks and may more than compensate for the energy cost required for the s-quark mass. This may (for high α_S) lead to a range of parameters where bulk strange quark matter is negatively charged. However, it has recently been shown [2] that even if negative strange quark matter is stable in bulk, and even if a negatively charged, low-mass strangelet were to be formed in a heavy-ion collision, that strangelet would not be able to eat up nuclei indefinitely and cause a disaster. This is because the same finite size effects that destabilize small strangelets relative to bulk matter [see (2) above] act via suppression of the s-quark wavefunction near the strangelet surface. This reduces the abundance of negatively charged s-quarks to a degree that makes the total strangelet charge positive for a very wide range of baryon numbers. A negative strangelet would simply stop growing by nucleus consumption when its mass reached the positively charged mass-regime, and no disasters would happen. The calculations behind this argument were performed within the MIT bag model, but the argument is likely to be much more generally valid because it simply relies on the quantum mechanical effect that massive particle wavefunctions are suppressed at a bag boundary.

(5) Recent developments in strong interaction theory indicate that interactions between quarks of different flavour and colour may lead to formation of quark Cooper pairs, resulting in a 'colour–flavour-locked' state that is more bound than 'ordinary' strange quark matter [6]. The arguments given in (1)–(3) above still apply if colour–flavour locking (CFL) is active, but the overall binding energy of strange quark matter will increase and the formation of strangelets will be more likely [7]. Bulk colour–flavour-locked strange quark matter is electrically neutral because the pairing energy is maximized when all three quark flavours have the same Fermi momentum, and therefore the same number density [8]. But for finite systems the net charge will be positive for the same reasons described in (4) above [7]. Thus these systems would not prove dangerous. CFL is apparently an inevitable consequence of QCD in the limit of infinite chemical potential. For quark chemical potentials of 300–400 MeV of relevance to strangelets it is not yet clear whether strange quark matter is ordinary or colour–flavour locked.

2.2 Constraints from cosmic rays and astrophysics

In principle, astrophysical observations may constrain the potential disaster scenarios for an accelerator like the LHC because similar conditions may occur in space. We shall review here the relevant arguments.

2.2.1 Cosmic rays

Cosmic-ray processes reach the energies and energy densities that will be encountered at the LHC and, therefore, they may provide limits on possible disaster scenarios. Such limits have been discussed in Refs. [1] and [3] and much of the analysis applies also to the LHC. Recent results obtained with a detector adding time-of-flight information to an array large enough to reach energies at and above the knee [9], approaching the LHC-equivalent energy region, confirm with improved accuracy that heavy ions have started to dominate the spectrum. Although the precise chemical composition is not known, the average value of A corresponds to that of magnesium, with ions at least as heavy as iron forming a substantial part. We summarize briefly here the main conclusions, taking into account the recent data from RHIC.

The first argument is based on the observation that high-energy cosmic rays collide often with the surface of astrophysical objects such as the moon. The total number of such collisions on the moon is huge compared to what is expected at the LHC and, since the moon is still there and made out of ordinary matter, one may try to use its existence to rule out the possibility of any dangerous process at the LHC. The argument is quite convincing, but one can still evade it in several ways.

First, the kinematic conditions are not identical. A complex hadronic system, like a strangelet, moving rapidly through matter, will not survive collisions with nuclei, so only strangelets nearly at rest relative to surrounding matter are potentially dangerous. At the LHC, they are the ones produced with almost zero rapidity while on the moon they are the ones in the target fragmentation region. Therefore, to compare the two, one needs some assumption on the rapidity distribution of strangelet production. For known hadrons these distributions fall like a power near the limiting values of the rapidity, so the authors of Ref. [1] assumed a distribution of the form

$$\frac{\mathrm{d}\Pi}{\mathrm{d}y} = Npy^a e^{-by} \ , \tag{1}$$

where a and b are parameters and N is a normalization constant chosen so that p gives half the total strangelet production probability per collision (the other half comes from the other half of the rapidity interval). The 'worst-case scenario' was assumed by the authors of Ref. [3] with the extreme case of a purely central production:

$$\frac{\mathrm{d}\Pi}{\mathrm{d}y} = p\delta(y - Y/2) \ , \tag{2}$$

where Y is the total rapidity interval. It is obvious that in this case the moon data put no limits on any collider. We want to emphasize, however, that we consider such a distribution to be totally unrealistic, although it cannot be rigorously disproved.

The second difference between the kinematical conditions on the moon and at the LHC has to do with the nature of the colliding nuclei. Neither cosmic rays nor moon soil are rich in very heavy nuclei like gold and the bulk of the data concern iron–iron collisions. There is no plausible reason why gold will produce strangelets but iron not; nevertheless, if one insists on gold in the incoming cosmic rays, one has to pay a penalty of the order of 10^{-5}.

The third problem is energy. The authors of Ref. [1] have argued that strangelet production should be at least as likely at AGS energies as at RHIC or the LHC, so beyond a certain threshold the incoming

energy should not be an important factor. We share this opinion, although we cannot give a rigorous proof. Nevertheless, if one considers only cosmic rays hitting the surface of the moon with energies equal to or higher than that of the LHC, one must pay a penalty factor given by the $E^{-3.4}$ behaviour of the cosmic-ray spectrum.

We conclude that the moon data alone can exclude the possibility of creating dangerous objects behaving like normal hadronic systems, but strangelets with completely exotic properties cannot be ruled out.

The authors of Ref. [3] have considered a second source of data, namely that of the collisions of cosmic rays in space, where secondary collisions cannot destroy the strangelets. The bound of strangelet production and growth follows from the fact that they will be eventually swept up in star formation and lead to the subsequent destruction of the star as it is converted into strange matter. This would be detectable as a supernova-like event. Despite the extremely small cross-section for interstellar cosmic-ray collisions, the fact that they may take place in an extremely large volume allowed Dar et al. [3] to obtain significant bounds on strangelet production and growth. However, this bound does not apply if the produced strangelet is unstable under baryon emission with a lifetime shorter than the time it takes for it to be swept up into a protostellar nebula [1]. Provided the lifetime lies between 10^{-7} seconds and several million years, a dangerous strangelet produced at the LHC would have time to stop in matter, stabilize and grow, but strangelets produced in space would not lead to disastrous stellar processes at an observable rate. In Ref. [1] the lifetime for such a baryon emission decay has been estimated with the conclusion that the lightest and most readily produced strangelets would decay with a much shorter lifetime and thus would not create a problem at RHIC or the LHC. Only the more massive strangelets would pose a problem, but they are much harder to produce.

To summarize, under plausible assumptions the cosmic-ray data exclude the possibility of dangerous processes in heavy-ion colliders like RHIC or the LHC, but the worst-case scenario cannot be excluded based on these data alone.

2.2.2 Astrophysics

Astrophysical observations set independent bounds on the existence of dangerous strangelets via three routes of argumentation.

(1) **Proof that bulk strange quark matter (regardless of charge) cannot be stable at zero pressure.** Since absolute stability of strange quark matter inevitably leads to the conclusion that all pulsars (and in general all compact objects normally believed to be neutron stars) are strange stars rather than neutron stars, proving that some compact object *is* a neutron star rather than a quark star would rule out the hypothesis of absolutely stable strange quark matter.

Strange stars with masses significantly below 1 solar mass are mainly confined by the strong interactions, with gravity playing a less important role. However, for the compact object masses typically created in supernova explosions ($M \approx 1.4$ solar masses), gravity is dominant for strange stars as well as neutron stars and the global stellar properties are not very different. In spite of several tentative 'proofs' and 'disproofs' of the existence of strange stars in the literature, the present situation is inconclusive. Possible ways of distinguishing between them in the future involve detailed comparisons of pulsar or X-ray binary radii and masses; the glitch phenomenon (sudden speed-up in pulsar rotation, for a long time thought to be very difficult to reconcile with strange stars, but apparently explainable with CFL); stellar neutrino cooling; the behaviour of r-mode instabilities; or the evolution of the braking index, which describes the time evolution of pulsar spin-down, and which has been suggested to show a very specific signature if a quark–hadron separation front moves as a pulsar spins down. This can only happen if quark matter is metastable rather than absolutely stable, and might therefore disprove the stability hypothesis if observed (though other phase transitions might mimic the signature).

While some of the signatures mentioned are promising in principle, the situation has been significantly complicated by the concept of CFL, which permits a much richer internal structure in strange stars than hitherto believed. At present there is no clearcut evidence for or against the existence of strange stars, and therefore for or against strange quark matter stability.

(2) **Proof that negatively charged bulk strange quark matter is not stable at zero pressure.** Since only negatively charged strange quark matter is potentially dangerous, it would suffice to prove that there are pulsar ('neutron' star) properties that are inconsistent with them being negatively (quark) charged strange quark matter. The distinctive feature here is the existence of an ordinary matter crust floating on top of the electrostatic potential around a quark-positive strange star, which has an electron 'atmosphere' outside the quark phase boundary. A negatively charged strange star would instead have an atmosphere of positrons and would not be able to sustain a crust of ordinary material. Some radiation phenomena from compact objects are believed to be consistent with a bare quark matter surface (possible for both positive and negative charge), others are most consistent with a crust. However, the evidence is far from conclusive, and the CFL mechanism may allow structures resembling a solid crust also in negatively charged quark stars (though at least for $\alpha_S = 0$, colour–flavour-locked quark stars would have zero charge [8]).

(3) **Proof that the flux of strangelets resulting from the cosmological quark–hadron transition, from strange star collisions in binaries, or from direct cosmic-ray production would exceed some observational constraints.** This is a very indirect probe. Significant limits on cosmological relic strangelets exist from Big Bang nucleosynthesis calculations, stellar evolution calculations (where strangelets could influence nucleosynthesis rates), and direct measurement limits from several cosmic-ray detectors. But since there are many reasons to expect that primordial strangelets did not form and/or survive the early Universe anyway (regardless of strange matter stability) these limits do not rule out the strange matter hypothesis. Most promising may be upcoming strangelet searches with the Alpha Magnetic Spectrometer (AMS) at the International Space Station (currently scheduled for launch in May 2005). Estimates of the flux of relativistic strangelets from strange star binary collisions (collisions in our Galaxy would occur roughly every 10^4 years if binary pulsars contain strange stars) show that AMS should be able to either detect strangelets or severely constrain the existence of strange stars before the start of the LHC heavy-ion programme.

2.3 Accelerator searches

2.3.1 Existing evidence

There have been essentially no new developments since the RHIC report was written. The most comprehensive search was carried out by E864 at the Brookhaven AGS with $11.5\,A$ GeV gold beams on platinum and lead targets. The final results [10] for the mass range $5 \leq A \leq 100$ and proper lifetimes > 50 ns are quoted as 0 candidates in 3×10^{10} central collisions, where central is defined as the upper 10% of the geometrical cross-section. The experiment was sensitive to both positive and negative electric charge and low values of $|Z/A| \leq 0.3$ characteristic of strangelets. An analogous search has been done by NA52 at the CERN SPS with $158\,A$ GeV lead beams on lead targets [11]. Again, no strangelets have been found under similar constraints to those quoted for the AGS; the sensitivity reached at present is 0 candidates in 10^8 to 10^9 collisions for positively charged strangelets, and 0 candidates in 10^9 to 10^{10} collisions for negatively charged. Notice, however, that both cover only large lifetimes, and both sensitivities fail to reach the coalescence and thermal model estimates by many orders of magnitude. Experiments at RHIC only started in 2000. RHIC will look for strangelets but the sensitivity reached at lower energies is not easily achievable.

2.3.2 Production mechanisms

As stated in the RHIC report, two mechanisms have been proposed for strangelet production in high-energy heavy-ion collisions: coalescence and strangeness separation or 'distillation'. For both mechanisms, we restrict ourselves to the central rapidity region. We will discuss coalescence first.

(1) **Coalescence.** In the late stages of the collision, after hadronization has occurred, strange and non-strange baryons which are close to each other in coordinate and momentum space may fuse to form nuclei and (multiple strange) hypernuclei. If a strangelet with the same quantum numbers (A, S) is more stable than the hypernucleus, the latter could act as a doorway to the former. Assuming in the estimates below that the transition from a hadronic to a quark structure occurs with unit probability, one surely gets a strong *overestimate* of the strangelet production probability, i.e. a very safe *upper limit*.

The basic physics idea of coalescence to form a nucleus A when N nucleons occupy the same point in (coordinate) space can be expressed as

$$Y_A = E\frac{\mathrm{d}^3 N_A}{\mathrm{d}P_A^3} = B_A (Y_N)^A = B_A \left(E\frac{\mathrm{d}^3 N_\mathrm{n}}{\mathrm{d}p_N^3} \right)^A , \tag{3}$$

which relates the yield Y_A of nuclei A with momenta P to the primordial nucleon invariant yield Y_N at momenta $p = P/A$ (assuming protons and neutrons with similar momentum spectra). The 'coalescence factor' B_A is proportional to the probability of finding the A nucleons at the same point in space. One therefore expects B_A to scale with the interaction volume like $1/V^{A-1}$ as long as the collision region is larger than the intrinsic size of the produced nucleus. B_A also includes the influence of binding energy of the final cluster; AGS data [12] point to a slight increase with increasing binding energy. In our extrapolation, we ignore this effect of binding energy. A transparent way to illustrate coalescence is to define a (dimensionless) 'penalty' factor P_F for coalescing an additional nucleon onto an existing cluster:

$$P_F = \frac{Y_A}{Y_{A-1}} = \left(\frac{B_A}{B_{A-1}} \right)_0 \times \frac{V_0}{V} \times Y_N . \tag{4}$$

This can be used to scale from a given cluster (index 0), e.g. deuterons, at a given energy to any other situation. A dependence of P_F on beam energy, in particular, is expected through the nucleon yield Y_N and the interaction volume V; assuming a constant freeze-out density, the associated freeze-out volume is proportional to the total multiplicity density $\mathrm{d}N_\mathrm{tot}/\mathrm{d}y$ of the produced particles, which indeed strongly depends on energy.

Reasonably systematic data on B_A and P_F exist for the AGS [12], the SPS [13] and already for RHIC [14]. For nuclei, penalty factors range from 1/50 at the AGS to about 1/300 at the SPS, for antinuclei (only $\overline{\mathrm{d}}$, $^3\overline{\mathrm{He}}$) from $1/(2 \times 10^5)$ at the AGS and 1/3000 at the SPS to about 1/1500 at RHIC (all values refer to central production). Nuclei and antinuclei have roughly the same values of B_A. The large dynamical range of the penalty factors simply reflects the large variations in Y_N and V. Within the experimental errors, the complete range of values is surprisingly well described by the scaling behaviour discussed above.

Given this nearly quantitative description of coalescence on a phenomenological basis, one can extrapolate to RHIC and LHC energies assuming that the penalty factors hold all the way up to larger values of A. To be on the safe side (an upper limit), we also do not introduce an additional penalty factor due to strangeness (in the RHIC report a factor of 0.2 was included). In other words, we take P_F to be flavour-independent. Such an assumption is supported by the thermal models discussed below. Also, already at full SPS energies, the yield of strange baryons is found to be roughly comparable to that of nonstrange baryons, and this is confirmed by the first RHIC data.

Starting from nucleons (roughly 100 per collision), we can then simply estimate the number of nuclei A per collision from

$$N_A \approx 100 \times P_F^{A-1} \ . \tag{5}$$

For a sample nucleus of $A = 20$, representing a strangelet of $A = 20$, $Z = -1$ and $S = 22$ as used in the RHIC report [1], and using extrapolated penalty factors of 1/900 and 1/1000 for central production at RHIC and the LHC based on the empirical scaling, one obtains yields of about 10^{-54} and 10^{-55} per collision. The closeness of these values is to be expected, since central production of baryons at collider energies is limited to baryon–antibaryon pair production, scaling with charged multiplicity density (and thereby $Y_N \sim V$), while additional baryons from stopping dominate at lower energies, increasing P_F. Even though we do not include an extra penalty for strangeness, the values we get are lower than those quoted in the RHIC report: This is mainly due to our inclusion of the correct scaling with beam energy. The sensitivity to the input assumptions is best illustrated by varying $1/P_F = 1000$ within $\pm 40\%$, consistent with the relative bandwidth in P_F describing the uncertainties of the experimental data set at low A. One then obtains a range of $10^{-55^{-3}_{+4}}$ $A = 20$ nuclei per collision as our final estimate for coalescence at LHC energies. The total number of heavy-ion (such as lead) collisions expected during ten years of LHC operation is 10^{11}. We shall use 10^{12} as an upper bound. As can be seen from Eq. (5), the result depends crucially on the value of A. However the bounds become dangerous only for A substantially lower than 10, for which no estimation produces stable, negatively charged strangelets.

Thermal models. An important feature of heavy-ion data is that hadron production is remarkably well described in terms of statistical models. Such models accurately predict the particle composition in terms of two parameters: the temperature T and the baryon chemical potential μ. These parameters are usually interpreted as characteristic of matter in chemical equilibrium at freeze out. Typical values are $T \approx 140$ MeV, $\mu \approx 540$ MeV at the AGS; $T \approx 170$ MeV, $\mu \approx 270$ MeV at the SPS; $T \approx 170$ MeV, $\mu \approx 50$ MeV at RHIC [15]. These numbers indicate that, as the energy increases, the chemical potential goes to zero and the temperature goes to a constant. This 'asymptotic' regime is practically reached at RHIC. In this regime, one expects the hadron yields to be the same for particles and antiparticles, and also to be independent of flavour.

The statistical approach can also be applied to the production of complex nuclei. The penalty factor P_F in the production cross-section for nucleus A compared to $A - 1$ is then a simple ratio of Boltzmann weights:

$$P_F = \frac{Y_A}{Y_{A-1}} \propto e^{-(m_N \mp \mu)/T} \ , \tag{6}$$

where m_N is the mass of the nucleon and μ the baryon chemical potential ($-\mu$ for antibaryons). In the asymptotic regime, one expects the penalty factor to be the same for matter and antimatter (because $\mu \approx 0$), and to be roughly the same for normal and multi-strange nuclei. There should also not be much difference between RHIC and the LHC. All of these features are quite compatible with the coalescence picture. Numerically, however, the parameters T and μ quoted above for the hadronization stage lead to penalty factors which are too large compared to the data systematics for nuclei by a factor of 3–6 depending on the case. A consistent description of the nuclear production cross-sections appears to require smaller values of both T and μ [13], for T somewhere between the chemical freeze-out temperature of ~ 170 MeV and the thermal freeze-out temperature of ~ 100 MeV. Since the systematics of these nuclear freeze-out parameters are not yet under complete control, we will not follow up more quantitative extrapolations.

It is however safe to conclude that the thermal model description, which provides a simple basis for extrapolation, is at least qualitatively in agreement with coalescence. The main outcome of the

present discussion is therefore the fairly robust prediction that production of strangelets should not differ much between RHIC and the LHC.

(2) **Distillation mechanism.** The distillation or strangeness-separation mechanism has been proposed in Ref. [16] as a specific model for strangelet production. It assumes the formation of a baryon-rich quark–gluon plasma which gets enriched in strangeness as it cools down.

There is so far no experimental evidence supporting this mechanism. If a quark–gluon plasma is formed in a collision, it will supposedly appear both in the central rapidity region and in the fragmentation regions where most of the net baryon number sits and where therefore the distillation mechanism should be most efficient. But production of negatively charged strangelets in the fragmentation regions is ruled out by moon data, as argued in the RHIC report. In central regions the net baryon number is small, even smaller at the LHC than at RHIC, and can only be significant through fluctuations, reducing the probability for the occurrence of distillation of strangelets.

It is not clear how much of this mechanism would contribute to the total rate of strangelet production, if it does at all. Besides, a large part of it would presumably be already included in coalescence calculations. We note in particular that, if the quark–gluon plasma phase is already reached at RHIC, such a mechanism should lead to an increase of the yields of all composite objects, including ordinary nuclei. There is however no evidence yet for a substantial enhancement of the production of composites on top of usual coalescence. Further data on this issue will become available from RHIC during the coming years.

2.4 Conclusions

- The coalescence picture provides the most reliable tool to extrapolate the present data to the energy range of the LHC. All accumulated evidence suggests strongly that no major change in the production of strangelets should be expected in going from RHIC to the LHC.
- Irrespective of any bounds coming from astrophysics or cosmic-ray data, the only dangerous scenario involving strangelets requires that for all A — from very low values, such as $A = 6$, all the way to $A = \infty$ — the negatively charged strangelets have lower mass than the corresponding neutral or positively charged ones and are stable, or long-lived.

3 GRAVITATIONAL EFFECTS

The RHIC report includes a discussion of whether heavy-ion collisions can assemble a sufficient concentration of matter to produce a black hole or other macroscopic object capable of growing due to classical gravitational forces. The authors conclude that such gravitational effects are totally negligible, being suppressed by inverse powers of the Planck mass $M_{\rm P}$. Recently there have been suggestions that the Planck mass is not a fundamental quantity but is derived from an underlying theory with more than four space-time dimensions. In such theories the higher dimensional Planck mass may be much smaller, raising the question of whether gravitational instabilities may develop much more readily. Given this we have re-examined the question whether the conditions will be such at the LHC as to produce stable black holes capable of accreting matter.

3.1 Black holes in four dimensions

We first review the situation in four dimensions. The general relativistic line element outside a spherical concentration of mass M is given in natural units by

$$\mathrm{d}s^2 = \mathrm{d}t^2(1 - \frac{2GM}{r}) - \frac{\mathrm{d}r^2}{\left(1 - \frac{2GM}{r}\right)} - r^2\mathrm{d}^2\Omega \; , \tag{7}$$

where G is Newton's constant. When $2GM/r = 1$ a horizon appears and the matter forms a black hole. This occurs at the Schwarzschild radius given by

$$R_{\rm S} = 2GM = \frac{M}{M_{\rm P}^2} \approx 10^{-33}{\rm cm}\left(\frac{M}{M_{\rm P}}\right) . \tag{8}$$

Heavy-ion collisions at the LHC will produce a concentration of energy over a length scale of $O(1\ {\rm TeV}^{-1}) \sim 10^{-17}$ cm. Thus we see that unless $M/M_{\rm P} > 10^{16}$ the energy will be distributed over a much larger volume than is needed to create a black hole. Since $M_{\rm P} \approx 10^{19}$ GeV we require

$$M > 10^{35}\ {\rm GeV}\ , \tag{9}$$

corresponding to 10^{32} nucleons with 1 TeV energy! Clearly the LHC or any conceivable accelerator will not produce collisions involving this number of nucleons.

Even this bound is weaker than is necessary because black holes decay (unless they carry a conserved quantum number Q, in which case extremal black holes with mass $M = Q$ are stable — see below for a discussion of this case). Let us estimate the decay lifetime. The temperature of a black hole is given by

$$T_{\rm BH} = \frac{M_{\rm P}^2}{M}\ . \tag{10}$$

Its thermal decay rate $\Gamma_{\rm D}$ is proportional to its area, giving

$$\Gamma_{\rm D} \approx T_{\rm BH}^4\, R_{\rm S}^2\ . \tag{11}$$

Unless the accretion rate is greater than this decay rate the black hole will decay harmlessly. An upper bound to the accretion rate is given by the energy density in the volume swept out by the black hole in one second. Assuming the limiting case where a black hole is moving with relativistic velocity the accretion rate is

$$\Gamma_{\rm A} \approx \pi R_{\rm S}^2 \rho\ , \tag{12}$$

where ρ is the mean density of the matter through which the black hole passes, which we take to be that of iron. Thus the condition for growth of the black hole, $\Gamma_{\rm A} > \Gamma_{\rm D}$, implies

$$\begin{aligned} T_{\rm BH} &< \rho^{1/4}\ , \\ M &> \frac{M_{\rm P}^2}{\rho^{1/4}}\ {\rm GeV} \approx 10^{42}\ {\rm GeV}\ . \end{aligned} \tag{13}$$

This is clearly a stronger bound than that of Eq. (9) and corresponds to the need to assemble 10^{39} nucleons with 1 TeV energy.

From this it is clear that classical gravitational effects are completely negligible for LHC energies and luminosities in the conventional theory of gravity. We turn now to a discussion of how these bounds change in the context of theories with large new space dimensions.

3.2 Large new space dimensions

It has recently been suggested [4, 17, 18] that there may be more than three spatial dimensions, provided the new space dimensions are compact with a size smaller than $R = 0.2$ mm. This latter bound comes from experiments looking for departures from Newton's law at short distances. In $(4 + d)$ space-time dimensions the gravitational potential of a mass m is given by

$$V(r) = \frac{KM}{M_*^{2+d}} \frac{1}{r^{1+d}}, \; r < R\,, \tag{14}$$

where M_* is the fundamental mass scale of the $(4+d)$ dimensional theory and $K = 16\pi/[(2+d)\Omega_{(2+d)}]$.

At distances larger than the compactification size of the new dimension, space-time effectively becomes four dimensional and the potential has the usual $1/r$ behaviour:

$$V(r) = \frac{KM}{M_*^{2+d} R^d} \frac{1}{r} \equiv \frac{GM}{r}, \; r > R\,. \tag{15}$$

Identifying the last two terms shows that the usual four-dimensional form of gravity applies with the Planck mass given by

$$M_{\mathrm{P}}^2 = \frac{2}{K} M_*^{2+d} R^d\,. \tag{16}$$

The extreme case is to choose $M_* = 1$ TeV, i.e. close to the electroweak breaking scale to avoid the hierarchy problem [17, 18]. With this choice one must have $R \leq 0.7$ mm for $d \geq 2$ in order to reproduce the correct value for the Planck mass. Recent experiments have probed the gravitational force law at scales down to 0.1 mm and disfavour the possibility with $d = 2$ but allow higher d and/or larger M_*.

3.3 Black holes in $(4 + d)$ dimensions

In $(4 + d)$ dimensions the general relativistic line element is modified from the form of Eq. (7) by the replacement of the gravitational potential GM/r with the higher dimensional potential given in Eq. (14). The corresponding Schwarzchild radius is given by [19]

$$\begin{aligned} R_{\mathrm{S}} &= \frac{K'}{M_*}\left(\frac{M}{M_*}\right)^{\frac{1}{1+d}} \\ &\approx \mathrm{TeV}^{-1}\left(\frac{M}{M_*}\right)^{\frac{1}{1+d}}\,, \end{aligned}$$

where

$$K' = \left(\frac{8\Gamma(\frac{3+d}{2})}{(2+d)\pi^{\frac{1+d}{2}}}\right)^{\frac{1}{1+d}}\,. \tag{17}$$

Compared with the expectation that heavy-ion collisions at the LHC will produce a concentration of energy over a length scale 1 $\mathrm{TeV}^{-1} \sim 10^{-17}$ cm, we see that the $(4 + d)$ black hole *will* be produced

if M is not much larger than 1 TeV. This is confirmed by detailed estimates in Refs. [20, 21]. Thus it is important to recalculate the stability bounds of black holes in $(4 + d)$ dimensions.

The temperature of the black hole is given by

$$T_{\rm BH} = \left(\frac{M_*}{M}\right)^{\frac{1}{1+d}} M_* \ . \tag{18}$$

The decay rate in this case is

$$\begin{aligned} \Gamma_{\rm D} &\approx T_{\rm BH}^{4+d} R_{\rm S}^{2+d} \\ &= M_*^2 \left(\frac{M_*}{M}\right)^{\frac{2}{2+d}} \ . \end{aligned} \tag{19}$$

A more precise determination of the decay rate may be obtained by summing over all the final states produced in black-hole decay, using a black-body spectrum modified by the grey-body (tunnelling) factors [22, 23, 24]. This gives the same dependence on the masses as in Eq. (19) but multiplied by a factor which varies from 0.7 for $d = 0$ to 0.1 for $d = 6$ [21, 25]. Such small corrections are completely irrelevant to the final conclusion so we will use the simple form of Eq. (19).

Since normal matter lives in four dimensions[1] the accretion rate has the same form as given in Eq. (12), so we immediately obtain the bound for a stable black hole given by

$$\rho \geq T_{\rm BH}^{4+d} R_{\rm S}^{d} \simeq T_{\rm BH}^{4} \tag{20}$$

as before. Using Eq. (18) gives the bound

$$M > M_* \left(\frac{10^4 M_*}{\rm GeV}\right)^{1+d} \ . \tag{21}$$

For the extreme case we have $M_* = 1$ TeV, so

$$M > 10^3 \times 10^{7(1+d)} \ {\rm GeV} \ . \tag{22}$$

Even for the case $d = 2$ (already disfavoured by experiment), the bound $M > 10^{24}$ GeV corresponds to 10^{21} nucleons of 1 TeV, again clearly beyond any accelerator. Thus we conclude that black-hole production does not present a conceivable risk at the LHC due to the rapid decay of the black hole through thermal processes.

3.4 Stable black holes and monopoles

One might worry that the discussion presented above fails for the case of a new conserved quantum number Q stabilizing the black hole. However, this is not the case because only extremal black holes are stable; others will rapidly decay to extremal black holes in the manner just discussed. Extremal black holes have $M = Q$ and so can only grow provided there is the source of the absolutely conserved quantum number Q. We know that normal matter does not possess such a quantum number and so there

[1]This estimate is unchanged even if normal matter does propagate in extra dimensions at short distances. In this case the effect of the extra dimensions is that normal matter has Kaluza–Klein excitations. That no such excitations have been observed implies these new states are very massive, $\geq O(1$ TeV$)$, and are thus not part of normal nuclear matter.

is no source of matter capable of causing the extremal black hole to grow, even if the LHC energy is capable of producing the new charge and thus a new stable form of matter.

However, such a new form of stable matter could cause a problem in a different way. Consider a magnetic monopole carrying a magnetic charge. It is conceivable in a theory with large extra dimensions that such a state could have a mass comparable to M_* and thus could be produced at the LHC. Moreover, magnetic monopoles can catalyse proton decay. Can this be a problem? At each catalysis event energy is released by the decaying proton, causing the monopole to move. It is straightforward to estimate the number of protons that could be destroyed before the monopole escapes the Earth. Monopoles are expected to have a strong cross-section with normal matter. As a result the mean free path of a monopole moving through iron is given by

$$\lambda = \frac{1}{\sigma_{\text{strong}}\rho} \simeq 1 \text{ cm} . \tag{23}$$

In the course of scattering N^2 times the monopole moves the distance λN and thus the number of scatters it experiences before escaping the Earth is determined from the condition $\lambda N = R_{\text{Earth}}$, corresponding to $N = 10^9$. In each collision a nucleon is destroyed so the escaping monopole will destroy 10^{18} nucleons: negligibly small compared to the total number of nucleons. Given this, we do not think it necessary to estimate the production rate for such new states because, even if there is no suppression in the production at the LHC, they do not present any conceivable threat.

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